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Ultra High Energy Cosmic Rays and de Sitter Vacua

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Abstract

The production of ultra high-energy cosmic rays in de Sitter invariant vacuum states is considered. Assuming the present-day universe is asymptoting toward a future de Sitter phase, we argue the observed flux of cosmic rays places a bound on the parameter α that characterizes these de Sitter invariant vacuum states, generalizing earlier work of Starobinsky and Tkachev. If this bound is saturated, we obtain a new top-down scenario for the production of super-GZK cosmic rays. The observable predictions bear many similarities to the previously studied scenario where super-GZK events are produced by decay of galactic halo super-heavy dark matter particles.

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I. INTRODUCTION

It has been known for a long time that de Sitter space possesses nontrivial vacuum states that are invariant under the symmetries of the space [1, 2, 3, 4], which we call the α-vacua. Physical applications of these states have recently been explored in the context of inflation, where they can lead to potentially observable corrections to the spectrum of cosmic microwave background fluctuations [5, 6, 7, 8, 9, 10, 11, 12, 13]. If such states can be relevant during inflation, it is natural to ask whether such states can have other observable consequences today. An initial study along these lines was made in [14] where the contribution to the cosmic ray spectrum was considered.

The cosmic ray spectrum has a feature at around 5×10^{18}eV where the power-law spectrum flattens from $E^{-3.2}$ to $E^{-2.8}$ as E increases, which suggests a transition from a galactic component of conventional astrophysical origin, to a component of extra-galactic origin. Some recent reviews of theoretical and experimental prospects for the study of these ultra high energy cosmic rays may be found in [15, 16]. Above about 10^{20}eV, protons rapidly lose energy due to their interaction with the cosmic microwave background, leading to the GZK cutoff [17, 18]. A handful of events above this bound have been observed, and there is still come controversy over whether or not the cutoff has been observed [19, 20, 21]. Extremely high energy cosmic rays $\gtrsim 10^{20}$eV are difficult to explain using conventional physics because likely sources lie outside the 100Mpc range of 10^{20}eV protons. A wide variety of scenarios have been proposed to account for the super-GZK events, which break down into two main classes: bottom-up mechanisms where charged particles are accelerated in large scale magnetic fields, and top-down mechanisms where exotic particles/topological defects produce extremely high energy cosmic rays via decay.

The α-vacua provide us with a new top-down mechanism for the production of extremely high energy cosmic rays. As noted in [22], a comoving detector can detect transitions of arbitrarily large energies (which we assume are cutoff near the Planck scale). The exception is the Bunch-Davies vacuum, where a detector measures a thermal response with a temperature of order the Hubble scale. These very high energy transitions in a generic α-vacuum can then account for production of extremely high energy cosmic rays.

Starobinsky and Tkachev [14] argued that if the α-vacua [43] do contribute to the ultra high energy cosmic ray spectrum at around 10^{20}eV, then the value of α becomes so tightly

constrained that it would not produce observable effects during inflation. In the present paper, we revisit this question and argue a much weaker constraint on α follows from cosmic ray observations. Our analysis also has bearing on the general question of what a low-energy observer will see in an α-vacua. The upshot of our analysis is that because production of cosmic ray flux necessarily violates de Sitter invariance, the production rate will be proportional to the background number density of matter, which leads to a much suppressed production rate versus the estimates of [14]. This rate is calculated in detail in section II.

From this result we infer bounds on α from cosmic ray observations. Assuming these bounds are saturated, we find the α-vacua give predictions very similar to extremely high energy cosmic ray production via decaying super-heavy dark matter in the galactic halo. This scenario has already been much studied in the literature [23, 24, 25]. We check that observable signals are out of reach in current neutrino/proton decay detectors. Finally we argue CPT violation in an α vacuum does not give rise to baryogenesis.

II. COMOVING DETECTOR IN DE SITTER SPACE

In a de Sitter invariant vacuum state, all correlators are invariant under the continuously connected symmetries of de Sitter space. In particular, this implies that $\langle n^\mu \rangle = 0$ for all 4-vectors n^μ, such as the number flux. Equivalently, the stress energy tensor in the de Sitter vacuum is proportional to the metric $T_{\mu\nu} \propto g_{\mu\nu}$. Since the metric is diagonal in comoving coordinates, this implies the absence of fluxes of energy or momentum. However a comoving detector nevertheless makes transitions due to its passage through the background spacetime, via the Unruh effect. We will model the injection spectrum of ultra high energy cosmic rays by viewing the universe today as de Sitter with $H = H_0$, the value of the Hubble parameter today. We treat the background density of ordinary Standard model matter as a small perturbation that explicitly breaks the de Sitter symmetry. Under certain circumstances, we can then treat these matter particles as Unruh detectors, which make transitions to highly excited states via interaction with the nontrivial vacuum state.

The α parameter in principle can depend on the species of field [26], which introduces a high degree of model dependence in the predictions [44]. For simplicity let us model the fields of observable matter by a single scalar field χ and assume that a different field ϕ (for example, the inflaton) is in a nontrivial α-vacuum, with coupling $\chi^2\phi$. We assume an order

1 coupling of ϕ to observable Standard model matter fields. The linear coupling of ϕ then allows us to treat the χ particles as an Unruh detectors (see [27] for a review).

As shown in [22] the rate at which an Unruh detector makes transitions is

$$\Gamma = N_\alpha^2 \left|1 + e^{\alpha + \pi\Delta E/H}\right|^2 \Gamma_0 \tag{1}$$

where Γ_0 is the result in the standard Bunch-Davies vacuum, $N_\alpha^2 = \frac{1}{1-e^{\alpha+\bar{\alpha}}}$. Γ_0 is Boltzmann suppressed by a $e^{-2\pi\Delta E/H}$ factor, so for large ΔE, Γ is proportional to $N_\alpha^2 \left|e^\alpha\right|^2$ times a power of ΔE. The injection spectrum is dominated by $\Delta E \sim M_c$ the field theory cutoff scale, which we have in mind to be of order the GUT scale 10^{16}GeV. When we integrate over ΔE, dimensional analysis then implies the total transition rate is

$$\Gamma \approx N_\alpha^2 \left|e^\alpha\right|^2 M_c. \tag{2}$$

Here we have assumed e^α is not so small that the 1 dominates in the $1 + e^{\alpha+\pi\Delta E/H}$ factor of (1). It is straightforward to generalize this expression to models with different couplings of α-vacuum species to observable matter, using the general formula (1).

III. ULTRA HIGH ENERGY COSMIC RAY PRODUCTION

Let us begin by reviewing the α-vacuum scenario, as described in [12]. During inflation, trans-Planckian effects [28] can lead to a de Sitter invariant state that differs from the conventional Bunch-Davies vacuum. If we invoke "locally Lorentzian" boundary conditions on modes, as described in [9, 11], one finds

$$e^\alpha \sim H/M_c. \tag{3}$$

This modifies inflationary predictions for the cosmic microwave background spectrum [9, 11, 12]. At the end of inflation, the value of the cosmological constant changes drastically. The squeezed state corresponding to the α-vacuum will then generate particles, producing an energy density of order [12, 14]

$$\varepsilon \sim N_\alpha^2 \left|e^\alpha\right|^2 M_c^4. \tag{4}$$

Provided $M_c \ll M_{Planck}$ this particle production does not overclose the universe, and instead can be thought of as some component of particle production during reheating. This energy

density will decay in a time of order $1/M_c$ (up to coupling dependent factors), as is typical of unstable particle production during reheating.

At much later epochs, it is still possible to have a residual α-vacuum present. If the universe asymptotes to a de Sitter universe with cosmological constant determined by the present value of H, the arguments of [9, 11] again apply and we can expect α given by (3) [10, 12]. One can then ask what phenomena observers today will see to indicate the presence of the α-vacuum. In [14] the assumption was made that (4) will be present for all times, and they used this to constrain α by matching with observed ultra high energy cosmic ray production. This assumption is equivalent to computing the energy density of the α-vacuum with respect to the Bunch-Davies vacuum, but gives the wrong result if the future asymptotic vacuum state is the α-vacuum. In this case, as we described in section II, no additional particle creation will be present in the de Sitter phase, and instead the particle production will be determined by (2), where we treat background matter as individual Unruh detectors.

In reality, the present universe is far from a pure de Sitter phase. The pure de Sitter estimate of the production rate nevertheless should be a reasonable order of magnitude estimate of the present rate of high energy particle production. Of course without a more detailed model for the dynamics that governs α we cannot make more precise statements.

Let us proceed then to compute the rate of high energy particle production in an α-vacuum. By the arguments of section II, we can then treat each Standard model particle as an Unruh detector, so (2) gives the rate of production per unit volume as

$$\frac{dn}{dt} = \Gamma n \tag{5}$$

where n is the number density of Standard Model particles [45]. In the situation of interest here, this density will be of order the baryon number density n_B which is typically

$$\begin{aligned} n_B &= 10\,\mathrm{m}^{-3} \approx H^2 M_{Planck}^2/m_p && \text{critical density} \\ &= 10^6\,\mathrm{m}^{-3} && \text{interstellar space} \end{aligned}$$

where m_p is the proton mass. Plugging in numbers, we find the dominant source of high energy cosmic rays will come from within our own galaxy due to interaction of visible and dark matter with the α vacuum. Many of the predictions will therefore be similar to the class of top-down models for ultra high energy cosmic ray (UHECR) production from decaying

super-heavy dark matter particles in the galactic halo. For a galactic halo of size r_{halo} , we find the flux received on earth will be of order

$$j \approx \Gamma n_B r_{halo}.$$

The experimental bounds coming from UHECR production gives $j\,E^2 \approx 10^{24}$ eV^2 m^{-2} s^{-1}sr^{-1} at $E \approx 10^{20}$ eV. Assuming $r_{halo} \approx 10^5$ light years, this translates into a bound

$$|e^{\alpha}| \lesssim 10^{-42}\left(\frac{10^{16}\mathrm{GeV}}{M_c}\right)^{1/2}. \tag{6}$$

This is to be compared with the "natural value" $e^{\alpha} \sim H/M_c \approx 10^{-61} M_{Planck}/M_c$ which is much smaller. We conclude then if ultra high energy cosmic ray production is to be accounted for by the α vacuum then the value of α must be much larger than its natural value.

It is interesting to ask if such a large value for α today might have other observable consequences. Let us also estimate the time needed for a neutrino style detector to see a nontrivial interaction with the α-vacuum. The interaction rate per baryon is (2) (taking $e^{\alpha} = H/M_c$, and $M_c = 10^{16}$GeV)

$$\Gamma = 10^{-76} s^{-1}$$

which is about 36 orders of magnitude smaller than current bounds on proton decay rate. For α saturating the bound (6) and $M_c = 10^{16}$GeV, we instead get

$$\Gamma = 10^{-44} s^{-1}$$

which is only 4 orders of magnitude smaller than the bounds on proton decay. We conclude that even if α is so large as to account for UHECR production, other means of direct detection will be difficult.

Finally, one might ask whether the present analysis has some impact on the spectrum of primordial inflation fluctuations. During the inflationary phase, the effect of the α-vacua on the primordial spectrum has been discussed in [7, 9, 11, 12], where it was found the amplitude of the spectrum was modulated by a factor of the form $1 + \mathcal{O}(H/M_c)$. The particle production effects described in the present paper will be absent in empty de Sitter, and we expect the effect will be a small correction to the energy density (4) in the context of slow-roll inflation. Note we already constrain (4) to be less than the vacuum energy density

during inflation [12]. Therefore we expect the particle production effects described here will have negligible impact on the spectrum of primordial fluctuations.

IV. BARYOGENESIS

An interesting feature of the α vacua is that they violate CPT symmetry when α is not a real number [3, 22]. This opens the possibility that the α vacua could be used to explain baryogenesis. If the CPT violation gives rise to particle/anti-particle mass differences, then baryogenesis could occur in thermal equilibrium, and might be relevant during the reheating phase at the end of inflation.

Greenberg [29] has argued that particle/anti-particle mass differences are only possible in flat space, if one gives up locality. We can apply these general results in the short wavelength limit of α vacuum propagators. As shown in [26] the interacting propagators give rise to local commutators in α vacua. In this limit, de Sitter symmetry becomes local Lorentz symmetry, so Greenberg's result will carry over. We conclude that α vacua do not lead to this type of baryogenesis.

V. DISCUSSION AND CONCLUSIONS

We have explored some of the phenomenological consequences of the present universe being in an α-vacua. It seems the most promising way to directly detect a residual value for α today is via observations of ultra high energy cosmic rays. As we have mentioned, many of the predictions will be similar to production of ultra high energy cosmic rays via decaying super-heavy dark matter particles in the galactic halo [23, 24, 25]. See [16, 30, 31] for some recent results, and more extensive references. Let us briefly discuss some of the features and constraints on this production mechanism.

Galactic halo cosmic ray production avoids the GZK cutoff, because the absorption length of ultra high energy protons is of order 100 Mpc. The cosmic rays typically do not have time to scatter before they reach us, so the observed spectrum should reflect the fragmentation function of the primary decay. This has been computed using Monte Carlo calculation in [30] including effects of supersymmetry. Perhaps the main problem one encounters in matching this with observation is that the fragmentation functions suggest the fraction of

gamma rays versus protons is too high versus the experimental bound [32, 33, 34]. This bound should become much better established in the upcoming Pierre Auger Observatory [35]. It is possible gamma rays lose energy more efficiently than protons over the scales of interest, which would ameliorate this problem. Searches for ultra high energy neutrinos should provide a more robust test of this scenario.

The observed arrival directions of UHECRs exhibit a high degree of isotropy on large scales, but clustering on smaller scales. This can be consistent with a clumpy distribution of dark matter in the galactic halo. The α vacuum scenario predicts additional anisotropy if the coupling to visible and dark matter is comparable, but these couplings are not well-constrained given the current level of understanding.

The main goal of the present work was to investigate observational constraints on a residual value of α today. These constraints easily allow for the theoretically preferred value of $|e^{\alpha}| \sim H(t)/M_c$. It is fascinating the α vacua may also lead to a possible explanation of the spectrum of UHECRs.

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[1] N. A. Chernikov and E. A. Tagirov, Annales Poincare Phys. Theor. **A9**, 109 (1968).

[2] E. A. Tagirov, Ann. Phys. **76**, 561 (1973).

[3] B. Allen, Phys. Rev. **D32**, 3136 (1985).

[4] E. Mottola, Phys. Rev. **D31**, 754 (1985).

[5] R. Easther, B. R. Greene, W. H. Kinney, and G. Shiu, Phys. Rev. **D64**, 103502 (2001), hep-th/0104102.

[6] R. Easther, B. R. Greene, W. H. Kinney, and G. Shiu, Phys. Rev. **D67**, 063508 (2003), hep-th/0110226.

[7] R. Easther, B. R. Greene, W. H. Kinney, and G. Shiu, Phys. Rev. **D66**, 023518 (2002), hep-th/0204129.

[8] L. Bergstrom and U. H. Danielsson, JHEP **12**, 038 (2002), hep-th/0211006.

[9] U. H. Danielsson, Phys. Rev. **D66**, 023511 (2002), hep-th/0203198.

[10] U. H. Danielsson, JHEP **12**, 025 (2002), hep-th/0210058.

[11] U. H. Danielsson, JHEP **07**, 040 (2002), hep-th/0205227.

[12] K. Goldstein and D. A. Lowe, Phys. Rev. **D67**, 063502 (2003), hep-th/0208167.

[13] G. L. Alberghi, R. Casadio, and A. Tronconi (2003), gr-qc/0303035.

[14] A. A. Starobinsky and I. I. Tkachev, JETP Lett. **76**, 235 (2002), astro-ph/0207572.

[15] L. Anchordoqui, T. Paul, S. Reucroft, and J. Swain, Int. J. Mod. Phys. **A18**, 2229 (2003), hep-ph/0206072.

[16] P. Bhattacharjee and G. Sigl, Phys. Rept. **327**, 109 (2000), astro-ph/9811011.

[17] K. Greisen, Phys. Rev. Lett. **16**, 748 (1966).

[18] G. T. Zatsepin and V. A. Kuzmin, JETP Lett. **4**, 78 (1966).

[19] M. Takeda et al., Phys. Rev. Lett. **81**, 1163 (1998), astro-ph/9807193.

[20] N. Hayashida et al., Astrophys. J. **522**, 225 (1999), astro-ph/0008102.

[21] T. Abu-Zayyad et al. (High Resolution Fly's Eye) (2002), astro-ph/0208243.

[22] R. Bousso, A. Maloney, and A. Strominger, Phys. Rev. **D65**, 104039 (2002), hep-th/0112218.

[23] V. A. Kuzmin and V. A. Rubakov, Phys. Atom. Nucl. **61**, 1028 (1998), astro-ph/9709187.

[24] M. Birkel and S. Sarkar, Astropart. Phys. **9**, 297 (1998), hep-ph/9804285.

[25] V. Berezinsky, P. Blasi, and A. Vilenkin, Phys. Rev. **D58**, 103515 (1998).

[26] K. Goldstein and D. A. Lowe, accepted for publication, Nucl. Phys. B (2003), hep-th/0302050.

[27] N. D. Birrell and P. C. W. Davies (1982), cambridge, Uk: Univ. Pr. 340p.

[28] J. Martin and R. H. Brandenberger, Phys. Rev. **D63**, 123501 (2001), hep-th/0005209.

[29] O. W. Greenberg, Phys. Rev. Lett. **89**, 231602 (2002), hep-ph/0201258.

[30] C. Barbot and M. Drees (2002), hep-ph/0211406.

[31] G. Sigl, Ann. Phys. **303**, 117 (2003), astro-ph/0210049.

[32] M. Ave, J. A. Hinton, R. A. Vazquez, A. A. Watson, and E. Zas, Phys. Rev. Lett. **85**, 2244 (2000), astro-ph/0007386.

[33] M. Ave, J. A. Hinton, R. A. Vazquez, A. A. Watson, and E. Zas, Phys. Rev. **D65**, 063007 (2002), astro-ph/0110613.

[34] K. Shinozaki et al., Astrophys. J. **571**, L117 (2002).

[35] J. W. Cronin, Nucl. Phys. Proc. Suppl. **28B**, 213 (1992).

[36] N. Kaloper, M. Kleban, A. Lawrence, S. Shenker, and L. Susskind, JHEP **11**, 037 (2002), hep-th/0209231.

[37] T. Banks and L. Mannelli (2002), hep-th/0209113.

[38] M. B. Einhorn and F. Larsen (2002), hep-th/0209159.

[39] M. B. Einhorn and F. Larsen (2003), hep-th/0305056.

[40] H. Collins, R. Holman, and M. R. Martin (2003), hep-th/0306028.

[41] K. Goldstein and D. A. Lowe (2003), hep-th/0308135.

[42] M. K. Parikh, I. Savonije, and E. Verlinde, Phys. Rev. **D67**, 064005 (2003), hep-th/0209120.

[43] Starobinsky and Tkachev analyze the situation where the vacuum is a mode number independent Bogoliubov transformation relative to the Bunch-Davies vacuum. Although they did not refer to them as such, these are the α-vacuum states studied earlier in [1, 2, 3, 4].

[44] A number of works have pointed out potential problems with the α-vacuum scenario [36, 37, 38, 39, 40]. These criticisms have been addressed in [10, 26, 41]. In particular, in [41] we show with a particular definition of the α-vacua in interacting scalar field theory, the theory is well-defined in a fixed de Sitter background. However when this theory is coupled to conventional gravity, problems with locality arise. Possible loopholes to this argument include: strongly coupled physics at the cutoff scale that render the perturbative arguments of [41] invalid; quantization on elliptic de Sitter [42], where one gives up time orientability; or other exotic non-local trans-Planckian effects that again fall outside the analysis of [41]. In the present

paper we only use properties of the free two-point function, and we assume some loophole to the arguments of [41] is in effect.

[45] The earlier calculation of [14] obtained $dn/dt \sim |e^{\alpha}|^2 H M_c^3$ (converting to our notation). Our result (5) is suppressed by a factor of order $HM_{Planck}^2/m_p M_c^2$, where the reader should recall H is the Hubble parameter today, and we have substituted the critical density for n.